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THE $s \rightarrow d\gamma$ TRANSITION IN KAON AND HYPERON DECAYS*

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We survey the possibilities of detecting the $s \rightarrow d\gamma$ transition in kaon and hyperon radiative decays. In the more frequent decays, like $K^+ \rightarrow \pi^+\pi^0\gamma$, $K^+ \rightarrow \pi^+\ell^+\ell^-$, the short-distance contribution is obscured by various long-distance transitions. Among the hyperon radiative decays, $\Xi^- \rightarrow \Sigma^-\gamma$ and $\Omega^- \rightarrow \Xi^-\gamma$ are the leading candidates for providing a window to the short-distance $s \rightarrow d\gamma$. The long-distance $s \rightarrow d\gamma$ transitions are also analyzed; their magnitude depends on the size of the deviation from a sum-rule among couplings of vector mesons to photons. The measurement of $\Omega^- \rightarrow \Xi^-\gamma$ could provide information on the relative size of the short and long-distance contributions to the magnetic component of $s \rightarrow d\gamma$.

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1. Introduction

Flavour changing single quark $Q \rightarrow q + \gamma$ transitions, induced by loop diagrams, are a basic feature of the electroweak standard model¹. The calculation² of the amplitude of such transitions from the electroweak model needs to be complemented by the inclusion of QCD corrections³. It is of obvious interest to identify physical processes in which these transitions contribute significantly. Such processes would test the standard model at the one loop level, as well as the procedure of administering to it the QCD corrections and then possibly provide windows to physics beyond the standard model.

Although the $s \rightarrow d\gamma$ transition was the first to be investigated^{1,2,4,5} in detail, with the aim of relating it to physical radiative processes of kaons and hyperons, it is the $b \rightarrow s\gamma$ transition⁶ which has conquered the limelight during recent years. Since it was shown⁷ that the enhancement provided by the QCD corrections to the $b \rightarrow s\gamma$ transition brings the inclusive $B \rightarrow X_s\gamma$ and the exclusive $B \rightarrow K^*\gamma$ modes into the realm of possible detection, the radiative penguin decays of the b-quark have received considerable attention. This has resulted in the calculation of the gluonic corrections fully to the leading order⁸ and partially at the next-to-leading order⁹. The recent measurements by the CLEO collaboration of the inclusive rate¹⁰ $Br(B \rightarrow X_s\gamma) = (2.32 \pm 0.57 \pm 0.35) \times 10^{-4}$ and of the exclusive decay¹¹ $Br(B \rightarrow K^*\gamma) = (4.5 \pm 1.5 \pm 0.9) \times 10^{-5}$ are in agreement with the theoretical expectations. An analysis¹² of the lowest order prediction for the inclusive mode gives $Br(B \rightarrow X_s\gamma)^{th} = (2.8 \pm 0.8) \times 10^{-4}$, with the error due mainly to the uncertainty in the choice¹³ of the renormalization scale. The completion of the next-to-leading order calculation will reduce significantly the theoretical uncertainty.

It is our purpose to review here the present status of the $s \rightarrow d\gamma$ transition and to analyze the possibility of its detection. Before turning to the main topic, we remark that the contribution of the $c \rightarrow u\gamma$ transition in charm decays has been analyzed¹⁴ recently in detail. It turns out that the electroweak $c \rightarrow u\gamma$ QCD uncorrected transition has a minuscule branching ratio of about 10^{-17} (vs. 10^{-4} in the $b \rightarrow s\gamma$ case). The inclusion of gluonic corrections raises dramatically this figure to nearly 10^{-11} ; however, this is still extremely small and as a result the domain of weak radiative decays of charmed particles is expected to be dominated by various long distance contributions.

2. The Flavour-Changing $Q \rightarrow q\gamma$ Transition in the Standard Model

The penguin amplitude for the transition of heavy quark Q to a light quark q , with quarks Q, q on the mass-shell is given^{1,2} by

$$\begin{aligned} \Gamma_{\mu}^{(Q \rightarrow q\gamma)} = & \frac{eG_F}{4\pi^2\sqrt{2}} \sum_{\lambda} V_{\lambda Q}^* V_{\lambda q} \bar{u}(q) [F_{1,\lambda}(k^2)(k_{\mu}\not{k} \\ & - k^2\gamma_{\mu}) \frac{1-\gamma_5}{2} + F_{2,\lambda}(k^2)i\sigma_{\mu\nu}k^{\nu}(m_Q \frac{1+\gamma_5}{2} + m_q \frac{1-\gamma_5}{2})] u(Q). \end{aligned} \quad (1)$$

$F_1 = \sum_{\lambda} V_{\lambda Q}^* V_{\lambda q} F_{1,\lambda}$ and $F_2 = \sum_{\lambda} V_{\lambda Q}^* V_{\lambda q} F_{2,\lambda}$ are the charge radius and magnetic form factors respectively and V_{ab} are CKM matrices. They were firstly calculated in the electroweak model by Inami and Lin¹, $F_{1,\lambda}$ and $F_{2,\lambda}$ giving the contribution of each internal quark in the loop. For the $s \rightarrow d\gamma$, the summation is over the u,c,t quarks. The term proportional to m_q is small and usually neglected. The F_1 term does not contribute to decays to real photons, however, it is relevant in weak lepton-pair decays, like $K^+ \rightarrow \pi^+ \ell^+ \ell^-$ ^{15,16}, $K_L^o \rightarrow \pi^o \ell^+ \ell^-$ ^{17,18}; it could also contribute, in principle, to leptonic decays of hyperons, like $\Sigma^+ \rightarrow p \ell^+ \ell^-$ ¹⁹ or $\Omega^- \rightarrow \Xi^- \ell^+ \ell^-$ ²⁰.

The QCD corrections are calculated by using an operator product expansion combined with renormalization group techniques. The effective Hamiltonian has the form

$$H_{eff}^{\Delta S=1} = -\frac{4G_F}{\sqrt{2}} V_{CKM} \sum_i c_i(\mu) O_i(\mu). \quad (2)$$

c_i are Wilson coefficients and O_i are a set of renormalized local operators, generated by the electroweak interactions and QCD. Physical amplitudes should be independent of the renormalization scale μ . The relevant operators in our case are¹⁸ $O_{1,2}$, the four QCD-penguin operators O_{3-6} , the electro- and chromomagnetic penguins $O_{7\gamma}, O_{8G}$, the electroweak penguins O_{7-10} and the semi-leptonic operators O_{9V}, O_{10A} . We give here those employed here more frequently and we refer the reader to Ref. 18 for the rest:

$$O_1 = (\bar{u}_{\alpha} \gamma_{\mu} P_L s_{\beta})(\bar{d}_{\beta} \gamma^{\mu} P_L u_{\alpha}); O_2 = (\bar{u}_{\alpha} \gamma_{\mu} P_L s_{\alpha})(\bar{d}_{\beta} \gamma^{\mu} P_L u_{\beta}) \quad (3a)$$

$$O_{7\gamma} = \frac{e}{16\pi^2} m_s (\bar{d}_{\alpha} \sigma_{\mu\nu} P_R s_{\alpha}) F^{\mu\nu}; O_{8G} = \frac{g_s}{16\pi^2} m_s (\bar{d}_{\alpha} \sigma_{\mu\nu} T_{\alpha\beta}^a P_R s_{\beta}) G^{\mu\nu} \quad (3b)$$

$$O_{9V} = (\bar{d}_{\alpha} \gamma_{\mu} P_L s_{\alpha})(\bar{e} \gamma^{\mu} e); O_{9A} = (\bar{d}_{\alpha} \gamma_{\mu} P_L s_{\alpha})(\bar{e} \gamma^{\mu} \gamma_5 e) \quad (3c)$$

where P_R, P_L are projection operators.

The Wilson coefficients are first computed perturbatively at $\mu = M_W$ in zeroth order of QCD and then evolved down to the scale of the process $\mu = m_0$ using the renormalization group equations (RGE). The solution of RGE to leading logarithmic order is given in terms of effective coefficients $c_k^{eff}(\mu = m_0)$, so that one obtains $c_k^{eff}(\mu = m_0) = u_{k\ell} c_\ell(M_W)$.

The function F_1 is related to c_{9V}^{eff} and F_2 is related to $c_{7\gamma}^{eff}$. In the next section, we shall discuss explicit calculations of the contribution of the short-distance $s \rightarrow d\gamma$ to several kaon decays.

3. Short-Distance $s \rightarrow d\gamma$ Contribution in Kaon Decays

The decay which has been firstly considered in detail¹⁵ as a possible handle on the $s \rightarrow d\gamma$ transition, with a virtual photon, is $K^+ \rightarrow \pi^+ \ell^+ \ell^-$. The decay amplitude can be written as

$$M(K^+ \rightarrow \pi^+ \ell^+ \ell^-) = \frac{ie^2 G_F}{4\sqrt{2}\pi^2} s_1 c_1 c_3 A s \cdot (p_K + p_\pi)^\mu \frac{\bar{u}(p_-) \gamma_\mu v(p_+)}{s} \quad (4)$$

where $s = (p_- + p_+)^2 = (p_K - p_\pi)^2$, $s_1 = \sin \theta_1$, etc. (CKM angles). A contains short-distance (SD) and long-distance (LD) contributions. The LD includes (a) transitions $K \rightarrow S\pi$ followed by $S \rightarrow \gamma^*$, where S is a strange intermediate hadronic state, (b) contributions involving a “ $K - \pi$ ” transition with pole, nonpole and contact terms and (c) contributions from strong penguin diagrams. These three classes can be calculated now fairly reliably¹⁶ and are all of the order 1. A similar result obtains in a chiral perturbation approach²¹. Moreover, the experimental rate and spectrum²² of $K^+ \rightarrow \pi^+ e^+ e^-$ agrees well with the predictions of Refs. [16,21]. Turning to the short-distance part, $s \rightarrow d\ell^+ \ell^-$ contains contributions from the electromagnetic penguin $s \rightarrow d\gamma^*$, the Z^o penguin $s \rightarrow dZ^{o*}$ and the “W box” diagram. The latter two may become significant in view of the heaviness of the top quark. However, since $|V_{ts}^* V_{td}| / |V_{cs}^* V_{cd}| \simeq 2 \times 10^{-3}$, it turns out that the contributions from Z^o -penguin and W-box are only a few percent of the main short-distance part which is due to the c-quark in the electromagnetic penguin loop. One finds¹⁶ $A_{SD} = 0.10$, hence the SD contribution to the rate is negligible and one cannot hope to detect its effect in rate or spectrum. This is also the reason for choosing (4) for the amplitude representation, despite the fact that in principle there is an axial part from O_{9A} .

There is, however, the possibility of detecting the short-distance contribution in $K^+ \rightarrow \pi^+ \mu^+ \mu^-$ from the interference of the one-photon LD part with the axial vector part coming from the SD contribution of Z^o -penguin and W-box diagrams. This induces²³ in this decay a very

small parity-violating longitudinal muon polarization asymmetry $\Delta_{LR} = (\Gamma_R - \Gamma_L)/(\Gamma_R + \Gamma_L)$, where $\Gamma_R(\Gamma_L)$ is the rate of producing a right-(left)-handed μ^- .

The above discussion on $K^+ \rightarrow \pi^+ \ell^+ \ell^-$ indicates that the LD contributions prevent an easy access to the SD ones. To reach these, one must turn to very rare decays, of which three have received considerable theoretical attention¹⁸ in recent years and are hopefully within reach of experimental detection²⁴ in the not too distant future. The three decays are $K_L^o \rightarrow \pi^o e^+ e^-$, $K^+ \rightarrow \pi^+ \nu \bar{\nu}$, $K_L^o \rightarrow \pi^o \nu \bar{\nu}$. We shall not discuss these decays in detail here except for a few general remarks, and we refer the reader to Refs. [18,24] for reviews on these modes. The $K_L^o \rightarrow \pi^o e^+ e^-$ decay contains¹⁷ a direct CP-violating part due to short-distance diagrams, an indirect CP-violating part induced via $K_s \rightarrow \pi^o e^+ e^-$ and a CP-conserving contribution. The direct CP-violating contribution is governed by the electromagnetic penguin operator O_{7V} as well as by O_{7A} and the strong penguin operators $O_3 - O_6$. It is expected to occur at a branching ratio of $\sim 4 \times 10^{-12}$, however, the other two contributions are estimated at comparable figures²⁵. The decay $K^+ \rightarrow \pi^+ \nu \bar{\nu}$ is the classic one for detecting SD contributions from Z-penguin and W-box, in view of the smallness of LD ones. Presently, after the next-to-leading order QCD contributions have been considered²⁶, one expects¹⁸ $0.6 \times 10^{-10} \leq Br(K^+ \rightarrow \pi^+ \nu \bar{\nu}) \leq 1.5 \times 10^{-10}$. The $K_L^o \rightarrow \pi^o \nu \bar{\nu}$ is also dominated by short-distance diagrams²⁷ and proceeds via direct CP-violation. The theoretical analysis predicts¹⁸ $10^{-11} \leq Br(K_L^o \rightarrow \pi^o \nu \bar{\nu}) \leq 5 \times 10^{-11}$.

We now turn to the part of $s \rightarrow d\gamma$ which occurs only in decays to real photons, namely the magnetic transition F_2 of Eq. (1). Firstly, with the normalization of Eqs. (1), (2), $F_2 = 2c_7$. The coefficient c_7^{eff} has been calculated recently by using the expressions of Buras et al.¹², with their anomalous dimension matrix, which gives $c_7^{eff}(\mu = m_0)$ in terms of $c_2(M_W), c_7(M_W), c_8(M_W)$. The last two coefficients contribute very little, since they are multiplied by $V_{ts}^* V_{td}$. Using $\alpha_s(M_W) = 0.12$, $\alpha_s(m_c) = 0.3$, $\alpha_s(\mu = m_0) = 0.9$ one obtains²⁸ $F_2 = 0.12$ which we shall use here. Other calculations^{3,29,30} give respectively 0.16, 0.20 and 0.08 for F_2 , using slightly different values as input in the calculation. This gives an indication on the sensitivity of c_7^{eff} to the input parameters, which in the leading log approximation is of the order of at least 50%.

The most frequent kaon decay sensitive to F_2 is $K^+ \rightarrow \pi^+ \pi^o \gamma$ ^{30,31}. The general form of the amplitude for the direct decay is³¹

$$M(K^+ \rightarrow \pi^+ \pi^o \gamma) = A \epsilon^{\mu\nu\sigma\tau} p_\mu^{(+)} p_\nu^{(0)} k_\sigma \epsilon_\tau + B [(p^{(+)} \cdot \epsilon)(p^{(0)} \cdot k) - (p^{(0)} \cdot \epsilon)(p^{(+)} \cdot k)] . \quad (5)$$

It has a parity-conserving magnetic transition A and a parity-violating electric transition B . The former appears to be the dominant contribution and the SD $s \rightarrow d\gamma$ contributes to it via F_2 . The calculation of Ref. 30 shows that this contribution, again, is only a few percent of the LD contributions³¹, which reproduce the experimental findings. Hence, here again, the branching ratio or the γ -spectrum cannot be used to reach the $s \rightarrow d\gamma$ transition.

4. Dynamics of Hyperon Radiative Decays

Since the more frequent radiative kaon decays are not useful for investigating the SD $s \rightarrow d\gamma$ transition, we turn to the baryonic (s,d,u) sector, in particular, the hyperon radiative decays. These decays, despite intensive attention during nearly thirty years, are still plagued with notorious discrepancies between theoretical models and experiment (see Ref. 32 for a recent review). The general form of the amplitude for such decays is

$$(H \rightarrow H'\gamma) = ieG_F\bar{u}(p')\sigma_{\mu\nu}(A + B\gamma_5)\epsilon^\mu k^\nu u(p) \quad (6)$$

where H, H' are spin $\frac{1}{2}$ baryons. As such, it is an ideal place for investigating the F_2 transition of Eq. (1).

For our purpose here, we note that the seven weak radiative decays which can occur within the octet and decuplet, can be divided³³ from a dynamical point of view into two groups, the “pole decays” ($\Sigma^+ \rightarrow p\gamma$, $\Lambda \rightarrow n\gamma$, $\Xi^o \rightarrow \Sigma^o\gamma$, $\Xi^o \rightarrow \Lambda^o\gamma$) and the “non-pole decays” ($\Xi^- \rightarrow \Sigma^-\gamma$, $\Omega^- \rightarrow \Xi^-\gamma$, $\Omega^- \rightarrow \Xi^{*-}\gamma$). The contribution of the SD $s \rightarrow d\gamma$ to the “pole decays”, occurring with branching ratios of the order of 10^{-3} , is negligible^{34,35}. The second group of decays involves particles $\Omega^- (sss)$, $\Xi^- (ssd)$, $\Sigma^- (sdd)$ with no u-valent quark, hence there are no W-exchange diagrams to generate poles. These decays may proceed via two-hadron intermediate states^{34,36}; since they are expected to have fairly low branching ratios ($10^{-4} - 10^{-5}$), it has been suggested³⁵ that the contribution of the SD $s \rightarrow d\gamma$ in the rate or in the asymmetry of decay might be detectable in some of them, especially $\Omega^- \rightarrow \Xi^-\gamma$.

So far, there are two measurements of the Ξ^- -decay, giving³² $Br(\Xi^- \rightarrow \Sigma^-\gamma) = 1.27 \pm 0.23 \times 10^{-4}$. A calculation³⁶ using the intermediate state of $(\Lambda\pi)$ only, gives a branching ratio $Br^{(th)}(\Xi^- \rightarrow \Sigma^-\gamma) = (1.8 \pm 0.4) \times 10^{-4}$ and an asymmetry coefficient ${}^{(th)}\alpha(\Xi^- \rightarrow \Sigma^-\gamma) = -0.13 \pm 0.07$. The uncertainties are due mainly to the possible additional contribution of $s \rightarrow d\gamma$. On the other hand, for Ω^- decay there is only an upper limit³⁷ $Br(\Omega^- \rightarrow \Xi^-\gamma) < 4.6 \times 10^{-4}$.

5. Short-Distance and Long-Distance $s \rightarrow d\gamma$ Contributions in $\Omega^- \rightarrow \Xi^-\gamma$

In this section we concentrate on the $s \rightarrow d\gamma$ contribution to the “non-pole” group of hyperon decays, in particular to $\Omega^- \rightarrow \Xi^-\gamma$. In addition to the SD contribution^{28,29}, there is also²⁸ a long-distance contribution of $s \rightarrow d\gamma$, which will be analyzed here.

Firstly, we note that the $s \rightarrow d\gamma$ transition (whether short- or long-distance) cannot be the dominant transition in *all* hyperon radiative decays. Indeed, if one assumes⁵ that its strength can be derived from the observed $\Sigma^+ \rightarrow p\gamma$ rate, the other radiative decays are predicted to be much larger than observed.

The SD amplitude for $s \rightarrow d\gamma$ is given by the F_2 term in Eq. (1). For the LD one, we consider a vector meson dominance (VMD) approximation to compute $s \rightarrow dV \rightarrow d\gamma$ transitions, along the lines discussed recently³⁸ for $b \rightarrow s\gamma$. We remark that a long-distance $s \rightarrow d\gamma$ transition is also one of the various contributions in $K^+ \rightarrow \pi^+ e^+ e^-$, as considered long ago by Vainshtein et al.¹⁵ with a somewhat different technique. Using the operators O_1, O_2 with u and c quarks and factorization, one obtains³⁸ for the amplitude $s \rightarrow d + V$

$$A(s \rightarrow d + V_i) = ig_{V_i}(q^2) \frac{G_F}{\sqrt{2}} a_2 V_{is} V_{id}^* \bar{d}\gamma_\mu (1 - \gamma_5) s \epsilon^\mu{}^+ \quad (7)$$

where a_2 is a QCD correction factor and $g_{V_i}(q^2)$ is defined $\langle V_i(a) | \bar{q}_i \gamma_\mu q_i | 0 \rangle = ig_{V_i}(q^2) \epsilon_\mu^+$. After extracting the transverse part of the amplitude in (7) to couple it to photons, the $s \rightarrow d\gamma$ amplitude including both SD and LD contributions is given by²⁸

$$\begin{aligned} A(s \rightarrow d\gamma) = A_{SD} + A_{LD} &= -\frac{eG_F}{\sqrt{2}} \bar{d}\sigma^{\mu\nu} \left[\left(\frac{m_s F_2}{8\pi^2} + \frac{va_2 C_{VMD} M_s}{M_s^2 - M_d^2} \right) P_R \right. \\ &\quad \left. + \left(\frac{m_d F_2}{8\pi^2} - \frac{va_2 C_{VMD} M_d}{m_s^2 - M_d^2} \right) P_L \right] s F_{\mu\nu} \end{aligned} \quad (8)$$

where $v = |V_{cs}^* V_{cd}^*| \approx 0.22$, $a_2 \simeq 0.5$, m_s , m_d are current masses and M_s , M_d are constituent quark masses.

$$C_{VMD} = \frac{2}{3} \sum_i \frac{g_{\psi_i}^2(0)}{m_{\psi_i}^2} - \frac{1}{2} \frac{g_\rho^2(0)}{m_\rho^2} - \frac{1}{6} \frac{g_\omega^2(0)}{m_\omega^2} \quad (9)$$

with the summation over the six narrow $c\bar{c}$ resonances. To calculate the physical decay $\Omega^- \rightarrow \Xi^-\gamma$ one uses²⁸ the amplitude (8) with SU(6) quark-model wave functions for the baryons and unit overlap. If we use only A_{SD} , with $F_2 = 0.12$, one finds $\Gamma_{SD}(\Omega^- \rightarrow \Xi^-\gamma) \simeq 6.4 \times 10^{-12}$ eV, far below the present experimental bound³⁷ of 3.7×10^{-9} eV. Using a value of $F_2 = 0.2$ and QCD sum rules to estimate the contribution of A_{SD} to the $\Omega^- \rightarrow \Xi^-\gamma$ decay, Nielsen et al.²⁹

find a value of Γ_{SD} which is about 8 times larger, but still two orders of magnitude below the experimental limit.

Using the full expression (8), the experimental upper limit determines $|C_{VMD}| < 0.01$ GeV². Assuming $g_V(0) \simeq g_V(m_V^2)$ for $V = \rho, \omega$, one obtains from (9) $\sum_i \frac{g_{\psi_i}^2(0)}{m_{\psi_i}^2} = 0.045 \pm 0.016$ GeV², about one sixth of the value at $m_{\psi_i}^2$. This confirms previous estimates summarized in Ref. 38 and shows that the right hand side of Eq. (9) is a remarkable sum rule holding to about 30%. This stems from the underlying input, which essentially include the GIM mechanism and SU(4)_F symmetry.

It is clear from the above discussion that the LD contribution to $s \rightarrow d\gamma$ may be significantly larger than the SD one and might even saturate the present experimental limit. It is important at this stage to mention additional possible contributions to the $\Omega^- \rightarrow \Xi^-\gamma$. Firstly, there is the unitarity limit, giving³⁴ $\text{Br}^{\text{unit}}(\Omega^- \rightarrow \Xi^-\gamma) > 0.8 \times 10^{-5}$. The dominant two-body intermediate state of $\Xi^0\pi^-$ contributes³⁴ $(1 - 1.5) \times 10^{-5}$ to the branching ratio and the strong penguin contribution is³⁹ about 4×10^{-6} or less. Hence, if the decay will turn out to be between the present upper limit of 4.6×10^{-4} and down to about 3×10^{-5} , the LD t-channel contribution of $s \rightarrow dV \rightarrow d\gamma$ is the dominant one. In this case the asymmetry of the decay should be⁴⁰ $\alpha(\Omega^- \rightarrow \Xi^-\gamma) = 0.4 \pm 0.1$, vs., say, the SD contribution which leads to an asymmetry 1. One should also add that the LD contribution of $s \rightarrow d\gamma$ with $|C_{VMD}| \leq 0.01$, is consistent with the observed hyperon decays²⁸.

In concluding, we stress that the more frequent kaon radiative decays are not profitable terrain for investigating the $s \rightarrow d\gamma$ transition and one must pursue the detection of the very rare decays, like $K^+ \rightarrow \pi^+\nu\bar{\nu}$, $K_L^0 \rightarrow \pi e^+e^-$, $K_L \rightarrow \pi^0\nu\bar{\nu}$ in order to check the SD/LD ratio in $s \rightarrow d\gamma$. On the other hand, the $\Omega^- \rightarrow \Xi^-\gamma$ decay might provide very useful information on the SD/LD ratio in $s \rightarrow d\gamma$. In fact, from (8) and the experimental upper limit on $\Omega^- \rightarrow \Xi^-\gamma$ we already know that $(\text{LD}/\text{SD}) \leq 25$ in this decay. Thus $s \rightarrow d\gamma$ represents an intermediate situation between the SD dominated $b \rightarrow s\gamma$, and $c \rightarrow u\gamma$ decays where the SD contribution is negligible.

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